

# STABILITY OF FINE TUNED HIERARCHIES IN STRONGLY COUPLED CHIRAL MODELS

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## ABSTRACT

A fine tuned hierarchy between a strongly coupled high energy compositeness scale and a much lower chiral symmetry breaking scale is a requisite ingredient in many models of dynamical electroweak symmetry breaking. Using a nonperturbative continuous Wilson renormalization group equation approach, we explore the stability of such a hierarchy against quantum fluctuations.

Many of the currently studied models of dynamical electroweak symmetry breaking include some strongly interacting sector acting at a high energy scale,  $\Lambda > 10TeV$ , which produces an essentially composite scalar bosonic degree of freedom. In addition, this dynamics is also supposed to play a non-trivial role in the electroweak symmetry breaking whose characteristic scale is much lower;  $\Lambda_F \simeq 250GeV$ . Thus these models require that a significant hierarchy can be established between these scales. Included in this class are strong extended technicolor models<sup>[1]</sup>, models of broken technicolor<sup>[2]</sup> and models involving heavy quark condensation<sup>[3-6]</sup>.

In general, the hierarchy is achieved by a fine tuning of parameters close to the critical value for the chiral symmetry breaking. A prototype of this behavior is exhibited by the Nambu Jona-Lasinio (NJL) model<sup>[7]</sup>, where a fine tuning of the four-fermion coupling allows the emergence of a chiral symmetry breaking scale far below the compositeness scale. A feature of this particular model which must be generic in any model which exhibits such a large hierarchy of scales is that the chiral symmetry phase transition be of second order. That is, in order for the hierarchy to be maintained and not have the electroweak scale driven to be of order  $\Lambda$ , it is necessary that the order parameter characterizing the chiral transition must remain zero as the theory is scaled from  $\Lambda$  into the infrared until one reaches the electroweak scale.

If, on the other hand, quantum fluctuations turn the transition first order at a scale  $e^{-t_0}\Lambda \gg \Lambda_F$ , then the order parameter will jump discontinuously to be of this value and it will be impossible to maintain the hierarchy all the way down to the electroweak scale. Instead the hierarchy will destabilize after  $t_0$  e-foldings. Such a situation is an example of the phenomenon studied by Coleman and Weinberg<sup>[8]</sup>. It is important to recognize that this question is distinct from that of the naturalness of the fine tuning of additive quadratic divergences. Even allowing for such fine tunings, this destabilization of hierarchies could prove problematic for many of the models of dynamical electroweak symmetry breaking. Thus, as emphasized by Chivukula, Golden and Simmons<sup>[9]</sup>, it becomes essential to explore when the hierarchy can be self consistently maintained and not destroyed by quantum fluctuations.

The minimal NJL model which is used in the minimal top condensate model<sup>[5]</sup> is known to exhibit a second order chiral transition. However, since a top quark mass of 174 GeV, as reported by the CDF collaboration<sup>[10]</sup>, is outside the range of the minimal model, it appears that if such models are to be phenomenologically viable, they must necessarily require extensions beyond the minimal version. Similarly, while the  $O(4)$  linear sigma model containing a single scalar quartic self coupling has been shown to exhibit a mean field second order chiral transition, the effective dynamics of strong extended technicolor models could require an effective Lagrangian containing more than one scalar self coupling. Thus we are led to investigate models containing multiple scalar quartic self couplings.

Following previous work <sup>[9,11,12]</sup>, we focus on a model possessing a global chiral  $U(2)_L \times U(2)_R$  symmetry which has two independent scalar quartic self

couplings. The model degrees of freedom include left and right handed chiral fermions  $\psi_{iL}$  and  $\psi_{iR}$ ,  $i = 1, 2$ , transforming as the fundamental  $(2, 0)$  and  $(0, 2)$  representations of the  $U(2)_L \times U(2)_R$  groups respectively which further carry the fundamental,  $N_C$ , representation of an asymptotically free gauged symmetry. Throughout our discussion, we shall neglect these gauge interactions which are considered to be very feeble at these high energy scales. We assume that, as a consequence of some unspecified dynamics acting at scale  $\Lambda$ , the  $U(2)_L \times U(2)_R$  chiral symmetry is spontaneously broken. This symmetry breaking is further assumed to produce a gauge singlet scalar composite  $\Sigma_{ij}$  which has the  $U(2)_L \times U(2)_R$  quantum numbers of the fermion bilinear  $\bar{\psi}_{jR}\psi_{iL}$ . That is  $\Sigma_{ij}$  transforms as the  $(\bar{2}, 2)$  under the chiral group. Its vacuum expectation value,  $\frac{v}{\sqrt{2}}\delta_{ij}$ , can be interpreted as an order parameter for the chiral symmetry breaking. Since we are assuming the chiral symmetry phase transition is second order, we are led to study a Ginzburg-Landau effective model whose Euclidean action at scale  $\Lambda$  is given by

$$S[\Sigma, \Sigma^\dagger, \psi, \bar{\psi}; 0] = \int d^4x \{ \text{tr}(\partial_\mu \Sigma^\dagger \partial^\mu \Sigma) + \bar{\psi} \gamma \cdot D \psi + V(x, y, 0) + \frac{\pi}{\sqrt{2}} g(0) (\bar{\psi}_L \Sigma \psi_R + \bar{\psi}_R \Sigma^\dagger \psi_L) \}, \quad (1)$$

where  $D_\mu \psi$  is the fermion gauge covariant derivative and the invariant potential function takes the form

$$V(x, y, 0) = \frac{1}{2} m^2(0) x + \frac{\pi^2}{12} \lambda_1(0) x^2 + \frac{\pi^2}{6} \lambda_2(0) y. \quad (2)$$

Here  $x = \text{tr}(\Sigma^\dagger \Sigma)$ ,  $y = \text{tr}(\Sigma^\dagger \Sigma)^2$  are the two independent  $U(2)_L \times U(2)_R$  invariants. A Coleman-Weinberg instability is signalled by the appearance of a non-trivial global minimum of the effective potential,  $V_{eff}(v)$ , with vanishing renormalized mass. If this occurs at the scale  $v = e^{-t_0} \Lambda$ , then the phase transition is driven first order by quantum fluctuations at this scale and one can technically achieve a hierarchy of only  $t_0$  e-foldings.

The conditions for the appearance of such a non-trivial global minimum of the effective potential with vanishing renormalized mass can be elegantly expressed in terms of the various renormalization group functions by employing a formalism due to Yamagishi<sup>[13]</sup>. In general, the (Yamagishi) stability function is defined as

$$Y(t_0) = -\frac{12}{\pi^2} e^{4t_0} \frac{dV_{eff}}{dt_0}, \quad (3)$$

with  $v = e^{-t_0}$ . The effective potential is minimized provided

$$Y(t_0) = 0 \quad ; \quad \frac{dY}{dt}|_{t=t_0} < 0 \quad (4)$$

and this minimum is a global one provided  $V_{eff}(v) < V_{eff}(0) = 0$ . Approximating the full effective potential by the 1-loop perturbation theory improved version which includes only the leading logarithmic radiative corrections of the dimension four operators appearing in the potential function of Eq. (2), while further holding the Yukawa coupling fixed, these Yamagishi conditions for the appearance of a global minimum at  $t_0$  reduce to

$$\begin{aligned} Y(t_0) &= [4\lambda_1 + 4\lambda_2 + \beta_1(\lambda) + \beta_2(\lambda)]|_{t=t_0} = 0 \\ [4 + \beta_1 \frac{\partial}{\partial \lambda_1} + \beta_2 \frac{\partial}{\partial \lambda_2}](\beta_1 + \beta_2)|_{t=t_0} &> 0 \quad ; \quad \lambda_2(t_0) > 0 \\ \lambda_1(t_0) + \lambda_2(t_0) &< 0, \end{aligned} \quad (5)$$

The first two lines of Eq. (5) are the fixed Yukawa coupling, 1-loop approximation version of Eq. (4) and guarantee a minimum at  $t_0$ , while the last inequality insures the minimum is a global one. Note that for other approximation schemes which are still restricted to include only the marginal, dimension four operators, the Yamagishi conditions of Eq. (5) are modified by the replacements  $\beta_i(t) \rightarrow \bar{\beta}_i(t) = \frac{\beta_i(t)}{1+\gamma(t)}$  and  $\lambda_i(t) \rightarrow \bar{\lambda}_i(t)$ , where  $\gamma(t)$  is the anomalous dimension of  $\Sigma$  and  $\bar{\beta}_i = \frac{\partial \bar{\lambda}_i}{\partial t}$ . Using the 1-loop perturbative  $\beta$  functions, the Yamagishi stability condition simply reduces to a curve in the  $\lambda_2(t)/g^2(0) - \lambda_1(t)/g^2(0)$  plane. The signal for the Coleman-Weinberg instability is that the renormalization group flow of the running couplings cross this stability curve. Using the 1-loop approximation with fixed Yukawa coupling, Chivukula et al.<sup>[9]</sup> found that, depending on the initial choice of couplings, the trajectories either

- (i) run to the infrared quasi fixed point near the origin,
- (ii) cross the stability line signaling a first order transition
- (iii) or simply run away in which case the model is ill defined.

For example, focusing on the region of coupling space corresponding to the initial couplings  $\lambda_2(0)$  and  $g^2(0)$  both large and positive and  $\lambda_1(0) = 0$ , the transition goes first order for sufficiently large  $\lambda_2(0)/g^2(0)$  ratio ( $\gtrsim 7$ ). Moreover, this occurs, in general, near to the compositeness scale. For example, choosing  $\lambda_2(0) = 10$  and  $g^2(0) = 1$  (and taking  $N_C = 3$ ), we display

Figure 1:  $Y(t)$  as a function of  $t$  for the initial parameters  $\lambda_1(0) = 0$ ,  $\lambda_2(0) = 10$ ,  $g^2(0) = 1$  computed using the 1-loop renormalization group improved effective potential with fixed Yukawa coupling.

in Fig. 1 the form of  $Y(t)$  as a function of  $t$ . This function is seen to cross the stability curve  $Y = 0$  at  $t_0 \sim 1.3$ . Thus the transition turns first order and the hierarchy destabilizes after  $\sim 1.3$  e-foldings (which corresponds to  $v \sim 0.27\Lambda$ ). However, since the couplings are very large, the 1-loop perturbative approximation can certainly be called into question. For instance, naively using the 2-loop perturbative renormalization group functions for these large initial couplings, the renormalization group trajectories do not cross the line but simply run away.

Clearly, some nonperturbative approximation scheme is required to properly deal with the system in the vicinity of the strong coupling compositeness scale. The purely bosonic  $U(2)_L \times U(2)_R$  model (no chiral fermions) has been simulated using lattice Monte Carlo techniques<sup>[11]</sup> and was seen to undergo a Coleman-Weinberg instability. An alternate approach which includes the chiral fermions has been advocated by Bardeen et al.<sup>[12]</sup>. Modelling the non-perturbative physics in the vicinity of the compositeness scale using a large  $N_C$  approximation, then at scale  $\Lambda$  the Yukawa coupling is seen to dominate. Retaining only it and the fine tuned scalar mass term needed to cancel the additive quadratic divergence, the model at scale  $\Lambda$  reduces to the minimal NJL model which is exactly soluble in the large  $N_C$  limit producing the run-

ning couplings  $\lambda_1(t) = 0$  and  $\lambda_2(t) = 3g^2(t) = \frac{48}{N_C t}$ . This model exhibits a (trivial) second order chiral transition and as such allows any sized hierarchy to be technically achieved. Note that for couplings consistent with the large  $N_C$  approximation, the 1-loop perturbative solution is driven to the infrared quasi fixed point and thus also does not cross the stability curve. As such it too allows for arbitrary large hierarchies to be tuned. Running the couplings using the large  $N_C$  approximation solution until the couplings have decreased sufficiently to be smoothly joined onto a perturbative running which also includes the running of the Yukawa coupling, it is found that the stability line is eventually crossed but only after a sizeable hierarchy ( $t_0 \sim 20 - 25$ ) has been established. Note that for this case of initially dominate Yukawa coupling (or for one of comparable size to the initial scalar self coupling), it is important to include the effects of its running which is to enhance the flow of the system toward the stability line. For smaller initial Yukawa couplings the effects of its running will not be nearly as important. Since the large  $N_C$  approximation is nonperturbative, the procedure of Bardeen et al. is a self consistent one. On the other hand, it can be reliably employed for only a very limited range of the initial parameter space.

An alternate nonperturbative method is provided by the continuous Wilson renormalization group equation (WRGE)<sup>[14–20]</sup> which has been extended to include chiral fermions<sup>[17]</sup>. The WRGE nonperturbatively relates the form of the Euclidean action at a scale  $e^{-t}\Lambda$  to the action at scale  $\Lambda$  for  $t > 0$ . It is derived by demanding that the physics, ie. correlation functions, remain unchanged as the degrees of freedom carrying momentum between scales  $\Lambda$  and  $e^{-t}\Lambda$  are integrated out. Thus either action can be used to equivalently describe the physics on all scales less than  $e^{-t}\Lambda$  and both actions lie on the same Wilson renormalization group trajectory. The lower scale action is constructed by appropriately changing the coefficients of the operators already present at scale  $\Lambda$ , as well as including new ones, in such a way so as to keep the physics unchanged. In general, the resultant action incorporates a complete set of local operators. This includes not only the relevant and marginal operators, but also irrelevant ones. The coefficient of each operator is fixed in terms of the initial action defined at scale  $\Lambda$ . The full WRGE is a very complicated integro functional differential equation and its analysis necessarily requires some simplifying approximations. One commonly used approximation is to work within a local action approximation<sup>[16–17]</sup> which ignores anomalous dimensions and derivative interactions while we further

neglect operators higher than bilinear in the fermion fields. In addition, in order to obtain a tractable analysis, we are forced to restrict attention to the case of a fixed Yukawa coupling. As such, the initial parameter space that we are able to investigate using this approach is restricted to be one where the initial scalar self coupling(s) dominates the initial Yukawa coupling which, however, can still be of order unity. Thus the non-perturbative WRGE we solve probes a region of initial parameter space intermediate between the lattice Monte Carlo simulations ( $g(0) = 0$ ) and the large  $N_C$  analysis ( $g(0)$  dominates) and our results can be viewed as complementary to those analyses. Taking into account the various simplifications, the action at scale  $e^{-t}\Lambda$  can be written as

$$S[\Sigma, \Sigma^\dagger, \psi, \bar{\psi}; t] = \int d^4x \{ \text{tr}(\partial_\mu \Sigma^\dagger \partial^\mu \Sigma) + \bar{\psi} \gamma \cdot D \psi + V(x, y, t) + \frac{\pi}{\sqrt{2}} g(0) (\bar{\psi}_L \Sigma \psi_R + \bar{\psi}_R \Sigma^\dagger \psi_L) \}, \quad (6)$$

where  $V(x, y, t)$  is the  $U(2)_L \times U(2)_R$  invariant potential function at this scale. The full WRGE<sup>[17]</sup> then reduces to a partial differential equation for this function of the form

$$\begin{aligned} \frac{\partial V}{\partial t} = & 4V - \Sigma_{ij} \frac{\partial V}{\partial \Sigma_{ij}} - \Sigma_{ij}^\dagger \frac{\partial V}{\partial \Sigma_{ij}^\dagger} \\ & + \frac{1}{16\pi^2} \text{tr} \ln(1 + W) - \frac{N_c}{4\pi^2} \text{tr} \ln(1 + \frac{\pi^2}{2} g^2(0) \Sigma^\dagger \Sigma), \end{aligned} \quad (7)$$

where  $W(\Sigma, \Sigma^\dagger)$  is the matrix of second derivatives

$$W = \begin{bmatrix} \frac{\partial^2 V}{\partial \Sigma_{ij}^\dagger \partial \Sigma_{kl}} & \frac{\partial^2 V}{\partial \Sigma_{ij}^\dagger \partial \Sigma_{kl}^\dagger} \\ \frac{\partial^2 V}{\partial \Sigma_{ij} \partial \Sigma_{kl}} & \frac{\partial^2 V}{\partial \Sigma_{ij} \partial \Sigma_{kl}^\dagger} \end{bmatrix} \quad (8)$$

The  $\text{tr} \ln$  terms containing  $W$  arise from integrating out the scalar modes, while those containing  $g^2(0)$  are due to integrating out the fermion modes. After a considerable amount of algebraic manipulation, the various determi-

nants can be explicitly evaluated yielding the WRGE<sup>[19]</sup>

$$\begin{aligned}
\frac{\partial V}{\partial t} = & 4V - 2xV_x - 4yV_y + \frac{1}{8\pi^2}\ell n[(1 + V_x + 2xV_y)^2 - 2(x^2 - y)V_y^2] \\
& + \frac{1}{16\pi^2}\ell n[(1 + V_x)^2 + 2x(1 + V_x)V_y + 2(x^2 - y)V_y^2] \\
& + \frac{1}{16\pi^2}\ell n[(1 + V_x)(1 + V_x + 6xV_y + 2xV_{xx} + 8yV_{xy} + 4x(3y - x^2)V_{yy}) \\
& \quad + 6(x^2 - y)V_y(3V_y + 2V_{xx} + 4xV_{xy} + 4yV_{yy}) \\
& \quad + 8(x^2 - y)(x^2 - 2y)(V_{xy}^2 - V_{xx}V_{yy})] \\
& - \frac{N_C}{4\pi^2}\ell n[1 + \frac{\pi^2}{2}xg^2(0) + \frac{\pi^4}{8}(x^2 - y)g^4(0)],
\end{aligned} \tag{9}$$

subject to the initial condition of Eq.(2). Here the subscripts denote differentiation with respect to that variable so that, for example,  $V_x = \frac{\partial V}{\partial x}$  etc. For  $t \geq 0$ , each action constructed using the  $V(x, y, t)$  satisfying this equation lies on the same Wilson renormalization group trajectory and produces the same physics on all scales less than  $e^{-t}\Lambda$ . Unfortunately, the solution to this equation is currently beyond our numerical abilities. Thus we make the further truncation of retaining terms only up to linear in  $y$  with coefficients which are arbitrary functions of  $x$ . Eq. (9) then reduces to two coupled equations which are of a similar (although considerably more complicated) form to what we previously solved in obtaining nonperturbative mass bounds<sup>[17,20]</sup>. While the truncations used are drastic and uncontrolled, they still include contributions from an infinite number of operators.

The resulting equations are then numerically solved for  $t$  values up to some  $t^*$ , where  $t^*$  lies in a region where  $V(x, y, t)$  is found to be linear in  $t - t^*$  with a slope of the same form as the linearized in  $t - t^*$  full 1-loop effective potential. The radiative corrections arising from the momentum modes less than  $e^{-t^*}\Lambda$  can then be satisfactorily incorporated using the 1-loop approximation to the effective potential. Note that this 1-loop effective potential also includes an infinite number of operators as is necessary to allow a smooth joining to the WRGE solution. So doing, the full effective potential is then constructed as

$$V_{eff}(x, y) = V(x, y, t^*) + V_{1-loop}(x, y, t^*), \tag{10}$$

where  $V_{1-loop}(x, y, t^*)$  is the analytically calculable 1-loop effective potential which accounts for the effects of the degrees of freedom carrying momentum



Figure 2:  $\lambda_1(t)$  as a function of  $t$  computed using the nonperturbative WRGE with fixed Yukawa coupling.

less than  $e^{-t^*} \Lambda$ . A Coleman-Weinberg instability is signalled by a non-trivial global minimum of  $V_{eff}(v) = V_{eff}(x, y)|_{y=\frac{1}{2}x^2=\frac{1}{2}v^4}$  with vanishing renormalized mass. The condition of vanishing renormalized mass,  $V_x^{eff}|_{x=y=0}$ , is achieved by an appropriate tuning of the parameter  $m^2(0)$  appearing in Eq. (2). If such a minimum appears at  $v = e^{-t_0} \Lambda$ , then the system can sustain a hierarchy only over  $t_0$  e-foldings.

Focusing on the specific initial couplings of  $\lambda_1(0) = 0$ ,  $\lambda_2(0) = 10$ ,  $g^2(0) = 1$ , we numerically integrated the WRGE and found that for  $t^* \sim 1.5$ ,  $V(x, y, t)$  was linear in  $t - t^*$  and smoothly joined onto the 1-loop effective potential. As an indication of this, we plot in Figs. 2-4 the behavior of the coefficients,  $\lambda_1(t) = \frac{6}{\pi^2} V_{xx}^{eff}|_{x=y=0}$ ,  $\lambda_2(t) = \frac{6}{\pi^2} V_y^{eff}|_{x=y=0}$  and  $\ell(t) = e^{-2t} (\frac{6}{\pi^2})^2 V_{xy}^{eff}|_{x=y=0}$ , of the two independent dimension four operators and one of the dimension six operators as a function of  $t$ . Note that we have included the explicit factor of  $e^{-2t} \sim \frac{1}{\Lambda^2}$  accompanying the dimension six operator in the definition of its coupling. As is clearly demonstrated, these exhibit a linear behavior in  $t - t^*$  for  $t^* \sim 1.5$ . The behavior of these couplings in this region is identical to that obtained using the linearized in  $t - t^*$  1-loop effective potential with anomalous dimensions neglected. Note that for small  $t$  values, the coefficient of the induced dimension six operator (and other irrelevant operators) are quite sizeable and play an important role in the dynamics. As  $t$  increases, the

Figure 3:  $\lambda_2(t)$  as a function of  $t$  computed using the nonperturbative WRGE with fixed Yukawa coupling.

effects of the irrelevant operators is diminished and the system is eventually attracted to the space spanned by the relevant and marginal operators only. Thus the contribution of the degrees of freedom with momentum less than  $e^{-t^*}\Lambda$  can then be included using the 1-loop approximation.

Since the WRGE approach includes an infinite number of operators, it would require an infinite dimensional space to plot the renormalization group flows. Thus we focus instead directly on the vacuum effective potential,  $V_{eff}(v)$ , defined in Eq. (10). In Fig. 5, we plot the generalized Yamagishi function  $Y(t) = -\frac{12}{\pi^2}e^{4t}\frac{dV_{eff}}{dt}$  as a function of  $t = -\ell n v$ . A first order transition is signaled if this function crosses the stability line  $Y = 0$ . It is seen that such a zero occurs at  $t_0 \sim 2.0$  and corresponds to a non-trivial global minimum of the effective potential at  $v \sim 0.14\Lambda$ . The value of  $t_0$  was seen to be insensitive to changes in the choice of  $t^*$ . Thus the transition goes first order and a hierarchy of only  $\sim 2.0$  e-foldings can be established. This result is in qualitative agreement with that found using the 1-loop perturbative approximation.

We have also investigated the phase structure of the model for other values of the initial coupling parameter space using the WRGE with fixed  $g^2(0)$ . For instance, for  $\lambda_1(0) = -2$ ,  $\lambda_2(0) = 10$ ,  $g^2(0) = 1$ , the Yamagishi function  $Y(t)$  crosses the stability line at  $t_0 = 1.9$  indicating a first order transition after

Figure 4:  $\ell(t)$  as a function of  $t$  computed using the nonperturbative WRGE with fixed Yukawa coupling.

Figure 5:  $Y(t)$  as a function of  $t$  for the initial parameters  $\lambda_1(0) = 0, \lambda_2(0) = 10, g^2(0) = 1$  computed using the nonperturbative WRGE with fixed Yukawa coupling.

1.9 e-foldings. Once again, this is in qualitative agreement with the 1-loop perturbative result of  $t_0 = 1.0$ . Similarly, for  $\lambda_1(0) = 2, \lambda_2(0) = 10, g^2(0) = 1$ , we find the chiral symmetry breaking transition turns first order at  $t_0 = 2.1$  while the 1-loop perturbative approximation gives  $t_0 = 1.7$ .

Using the nonperturbative WRGE, we have studied the stability of fine tuned hierarchies in a  $U(2)_L \times U(2)_R$  chiral model for a range of initial couplings satisfying  $\lambda_2(0) \gg g^2(0) \sim |\lambda_1(0)| \sim 1$ . In general, we found that the chiral symmetry phase transition turns first order in close proximity to the compositeness scale. This is in qualitative agreement with the results obtained using the 1-loop improved effective potential. In performing our analysis, we were forced to employ various truncations in order to render the numerical integrations tractable. Clearly the sensitivity of our results to these various approximations needs further scrutiny. Moreover, some of these truncations must be eliminated before we are able to probe the region of initial parameter space with larger Yukawa couplings. We hope to address these issues in future studies.

We thank Sergei Klebnikov for many enjoyable useful discussions and Bijan Haeri for assistance with the numerical computations. This work was supported in part by the U.S. Department of Energy under grant DE-AC02-76ER01428 (Task B).

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